Neutrinos in astroparticle physics

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Abstract. I briefly discuss the role of neutrinos as probes in astroparticle physics and review the status of neutrino oscillation parameters as of June 2006, including recent fluxes, and latest SNO, K2K and MINOS results. I comment on the origin of neutrino masses in seesaw-type and low-scale models and mention some of their experimental signals.

INTRODUCTION

Neutrinos play a central role as probes in astroparticle physics and are basic indicators of what may lie ahead of the Standard Model (SM). The discovery of neutrino oscillations comes mainly from the study of "heavenly" neutrinos [1, 2, 3], and has been brilliantly confirmed by laboratory data from reactors [4] and accelerators [5, 6].

Here I summarize the status of the interpretation of the current neutrino data within the simplest CP-conserving three-neutrino oscillation scenario. In addition to a determination of the solar and atmospheric oscillation parameters θ_{12} & $\Delta m^2_{\rm SOL}$ and θ_{23} & $\Delta m^2_{\rm ATM}$, one gets a constraint on the angle θ_{13}. Together with the small ratio $\Delta m^2_{\rm SOL}/\Delta m^2_{\rm ATM}$ the angle θ_{13} holds the key for future searches for CP violation in neutrino oscillation. The growing precision of oscillation experiments also opens good prospects for improved robustness tests, probing unitarity violation [7] and other forms of non-standard neutrino interactions.

The search for lepton number violating processes such as neutrinoless double beta decay [8, 9] ($0\nu\beta\beta$) constitutes a very important goal for the future, as this will probe whether neutrinos are Dirac or Majorana particles, irrespective of the mechanism that induces their mass. This is known as the "black-box" theorem [10]. In addition, $0\nu\beta\beta$ will be sensitive to the absolute scale of neutrino mass and to CP violation induced by the so-called Majorana phases [7], inaccessible in conventional oscillations [11, 12, 13].

STATUS OF NEUTRINO OSCILLATIONS

The discovery of oscillations marks a turning point in particle and nuclear physics and implies that neutrinos have masses. This possibility has been first suggested by theory in the early eighties, both on general grounds and on the basis of different versions of the seesaw mechanism [14, 15, 16, 7, 17, 18]. The basic ingredient is the lepton mixing

TABLE 1. Neutrino oscillation parameters as of June 2006, from Ref. [23].

parameter	best fit	3σ range
Δm^2_{21} [10^{-5} eV^2]	7.9	7.1–8.9
Δm^2_{31} [10^{-3} eV^2]	2.6	2.0–3.2
$\sin^2\theta_{12}$	0.30	0.24–0.40
$\sin^2\theta_{23}$	0.50	0.34–0.68
$\sin^2\theta_{13}$	0.00	≤ 0.040

matrix, whose simplest unitary 3-dimensional form is given as [7]

$$K = \omega_{23}\omega_{13}\omega_{12} \tag{1}$$

where each ω is effectively 2×2, characterized by an angle and a CP phase. Majorana phases do not affect oscillations and, moreover, current neutrino oscillation data have no sensitivity to the remaining Dirac CP violation phase. Thus we set the three phases to zero. In this approximation oscillations depend on the three mixing parameters $\sin^2\theta_{12}, \sin^2\theta_{23}, \sin^2\theta_{13}$ and on the two mass-squared splittings $\Delta m^2_{\mathrm{SOL}} \equiv \Delta m^2_{21} \equiv m^2_2 - m^2_1$ and $\Delta m^2_{\mathrm{ATM}} \equiv \Delta m^2_{31} \equiv m^2_3 - m^2_1$ characterizing solar and atmospheric neutrinos. The hierarchy $\Delta m^2_{\mathrm{SOL}} \ll \Delta m^2_{\mathrm{ATM}}$ implies that, to a good approximation, one can set $\Delta m^2_{\mathrm{SOL}} = 0$ in the analysis of atmospheric and accelerator data, and $\Delta m^2_{\mathrm{ATM}}$ to infinity in the analysis of solar and reactor data.

Interpreting the data requires good calculations of the corresponding fluxes [19, 20], neutrino cross sections and response functions, as well as an accurate description of neutrino propagation in the Sun and the Earth, taking into account matter effects [21, 22].

The resulting three–neutrino oscillation parameters obtained in the global analysis are summarized in Fig. 1. The analysis includes all new neutrino oscillation data, as of June 2006, as described in Appendix C of hep-ph/0405172 (v5) [23]. These include new Standard Solar Model [24], new SNO salt [25], latest K2K [5] and MINOS [6] data. In the upper panels of the figure the $\Delta\chi^2$ is shown as a function of the three mixing parameters $\sin^2\theta_{12}, \sin^2\theta_{23}$, $\sin^2\theta_{13}$ and two mass squared splittings $\Delta m^2_{21}, \Delta m^2_{31}$, minimized with respect to the undisplayed parameters. The lower panels show two-dimensional projections of the allowed regions in the five-dimensional parameter space. In addition to a confirmation of oscillations with $\Delta m^2_{\mathrm{ATM}}$, accelerator neutrinos provide a better determination of $\Delta m^2_{\mathrm{ATM}}$ as one can see by comparing dashed and solid lines in Fig. 1. Clearly MINOS [6] leads to an improved determination and a slight increase in $\Delta m^2_{\mathrm{ATM}}$. On the other hand reactors [4] have played a crucial role in selecting large-mixing-angle (LMA) oscillations [26] out of the previous "zoo" of solutions [27]. Table 1 summarizes the current best fit values and the allowed 3σ ranges that follow from the global fit.

Note that in a three–neutrino scheme CP violation disappears when two neutrinos become degenerate or when one of the angles vanishes [28]. As a result CP violation is doubly suppressed, first by $\alpha \equiv \Delta m^2_{\mathrm{SOL}}/\Delta m^2_{\mathrm{ATM}}$ and also by the small value of θ_{13}. The left panel in Fig. 2 gives the parameter α, as determined from the global χ^2 analysis.

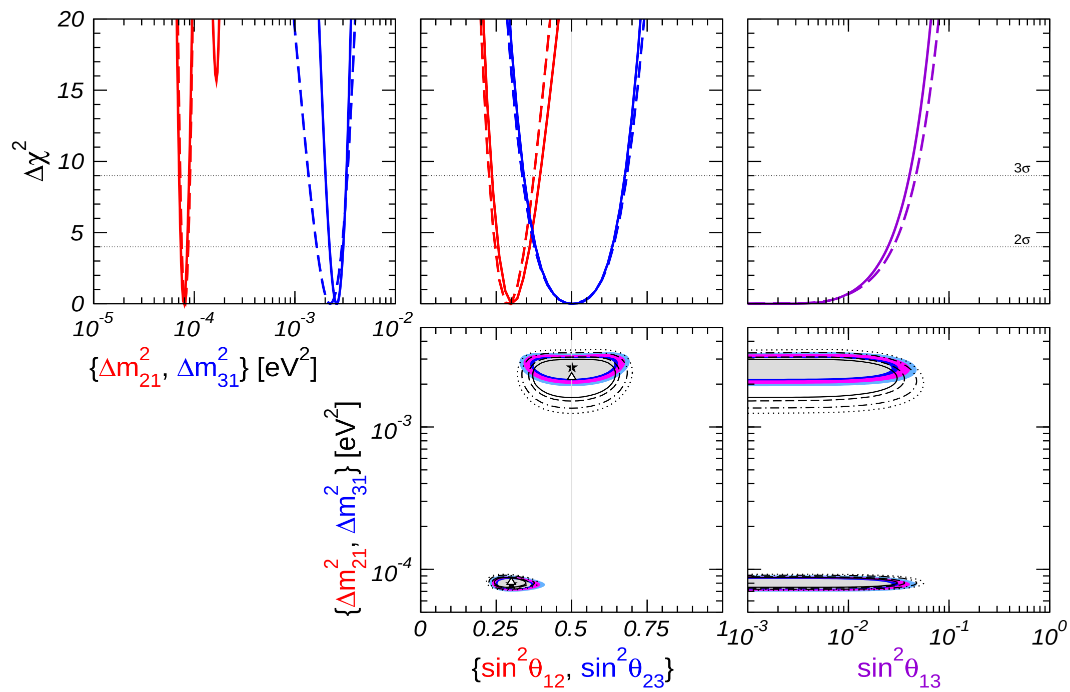

FIGURE 1. Current neutrino oscillation regions at 90%, 95%, 99%, and 3σ C.L. for 2 d.o.f. from Ref. [23]. In top panels $\Delta\chi^2$ is minimized with respect to undisplayed parameters.

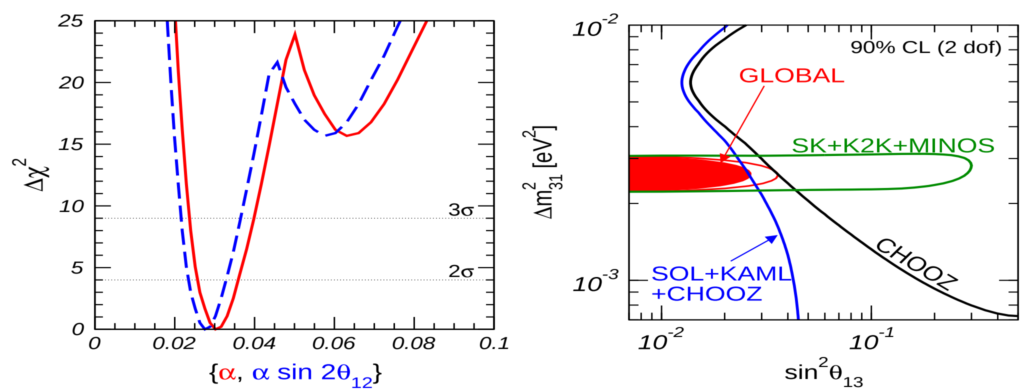

FIGURE 2. $\alpha \equiv \Delta m^2_{\rm SOL}/\Delta m^2_{\rm ATM}$ and $\sin^2\theta_{13}$ bound from the updated analysis given in Ref. [23].

The right panel shows the impact of different data samples on constraining θ_{13}. One sees that for larger $\Delta m^2_{\rm ATM}$ values the bound on $\sin^2\theta_{13}$ is dominated by CHOOZ, while for low $\Delta m^2_{\rm ATM}$ the solar and KamLAND data become quite relevant.

There is now an ambitious long-term effort towards probing CP violation in neutrino oscillations [29, 30, 31]. As a first step, upcoming reactor and accelerator long baseline experiments aim at improving the sensitivity on $\sin^2\theta_{13}$ [32]. An alternative possibility involving the day/night effect studies in large water Cerenkov solar neutrino experiments such as UNO, Hyper-K or LENA has also been suggested [33].

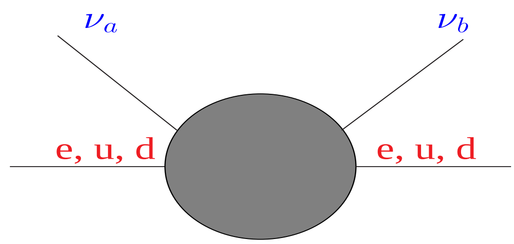

FIGURE 3. Non-standard neutrino interactions arise, e. g., from the non-unitary structure of charged current weak interactions characterizing seesaw-type schemes [7].

Reactor neutrino data have played a crucial role in testing the robustness of solar oscillations vis a vis astrophysical uncertainties, such as magnetic fields in the radiative [34, 35, 36] and convective zone [37, 38, 39], leading to stringent limits on neutrino magnetic transition moments [40]. KamLAND has also played a key role in identifying oscillations as "the" solution to the solar neutrino problem [26] and also in pinning down the LMA parameter region among previous wide range of oscillation solutions [27].

However, there still some fragility in the interpretation of the data if sub-weak strength ($\sim \varepsilon G_F$) non-standard neutrino interaction (NSI) operators (Fig. 3) are included. Indeed, most neutrino mass generation mechanisms imply the existence of such dimension-6 operators. They can be of two types: flavour-changing (FC) and non-universal (NU). Their presence leads to the possibility of resonant neutrino conversions even in the absence of neutrino masses [41]. While model-dependent, the expected NSI magnitudes may well fall within the range that will be tested in future precision studies [32]. For example, in the inverse seesaw model [42] the non-unitary piece of the lepton mixing matrix can be sizeable, hence the induced non-standard interactions. Relatively sizable NSI strengths may also be induced in supersymmetric unified models [43] and models with radiatively induced neutrino masses [44, 45].

The determination of atmospheric neutrino parameters $\Delta m^2_{\rm ATM}$ and $\sin^2\theta_{\rm ATM}$ is hardly affected by the presence of NSI on down-type quarks, at least within the 2–neutrino approximation [46]. Future neutrino factories will substantially improve this bound [47].

In contrast, the determination of solar neutrino parameters is not quite robust against the existence of NSI [48], even if reactor data are included. One can show that even a small residual non-standard interaction in this channel has dramatic consequences for the sensitivity to θ_{13} at a neutrino factory [49]. Improving the sensitivities on NSI constitutes at a near detector or via coherent neutrino scattering off nuclei [50] a window of opportunity for neutrino physics in the precision age.

THE ORIGIN OF NEUTRINO MASS

Here I briefly discuss the theory of neutrino mass and mention some recent attempts at predicting neutrino masses and mixing.

Light Dirac neutrinos

Gauge theories prefer Majorana neutrinos [7]. This statement holds irrespective the detailed model of neutrino mass generation. The emergence of Dirac neutrinos would constitute a surprise, indicating the existence of a fundamental lepton number symmetry whose origin should be understood. Without a specific reason, the appearance of such symmetry would be “accidental”.

Nevertheless there are interesting ideas for generating light Dirac neutrinos. For example, theories involving large extra dimensions offer a novel scenario to account for small neutrino masses [51, 52, 53, 54, 55, 56, 57]. According to this picture, right-handed neutrinos propagate in the bulk while left-handed neutrinos, being a part of the lepton doublet, live only on the SM branes. As a result, neutrinos can naturally get very small Dirac masses via mixing with a “bulk” fermion.

Light Majorana neutrinos

Charged fermions in the SM come in two chiral species to provide their mass after the electroweak symmetry breaks through the nonzero vacuum expectation value (vev) of the Higgs scalar doublet $\langle\Phi\rangle$. Neutrinos do not. There is, however, an effective lepton number violating dimension-five operator $\lambda L\Phi L\Phi$ in Fig. 4, which can be added to the SM (here L denotes any of the lepton doublets) [16]. After the Higgs mechanism this

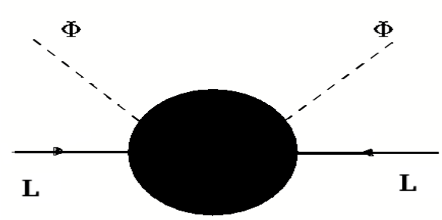

FIGURE 4. Dimension five operator responsible for neutrino mass [16].

induces Majorana neutrino masses $\propto \langle\Phi\rangle^2$, thus providing a natural way to account for the smallness of neutrino masses, irrespective of their specific origin. Little more can be said from first principles about the *mechanism* giving rise to this operator, its associated mass *scale* or its *flavour structure*. Its strength λ may be suppressed by a large scale M_X in the denominator (top-down) scenario, leading to $m_\nu = \lambda_0 \frac{\langle\Phi\rangle^2}{M_X}$, where λ_0 is some unknown dimensionless constant. Gravity has been argued to break global symmetries and thus could induce the dimension-five operator, with $M_X = M_P$, the Planck scale [58]. In this case the magnitude of the resulting Majorana neutrino masses are too small.

Alternatively, the strength λ may be suppressed by small parameters (e.g. scales, Yukawa couplings) and/or loop-factors (bottom-up scenario) with no need for a large scale. Both classes of scenarios have been reviewed in [59]. Here is a brief summary.

Seesaw-type models

The most popular top-down scenario is the seesaw [14]. The idea is to generate the dim-5 operator by the exchange of heavy states, either fermions (type-I) or scalars (type-II), typically both, as shown in Fig. 5. The main point is that, as the masses of the intermediate states go to infinity, neutrinos become light [15]. The seesaw provides a simple realization of Weinberg's dim-5 operator [16]. It can be implemented in many ways, with explicitly or spontaneously broken B-L, gauged or not; with different gauge groups and multiplet contents, minimal or not; with its basic scale large or small. Seesaw with gauged B-L broken at large scale is but one possibility. I have no space here for a detailed discussion, those interested in a short seesaw "Kamasutra" may consult Ref. [59].

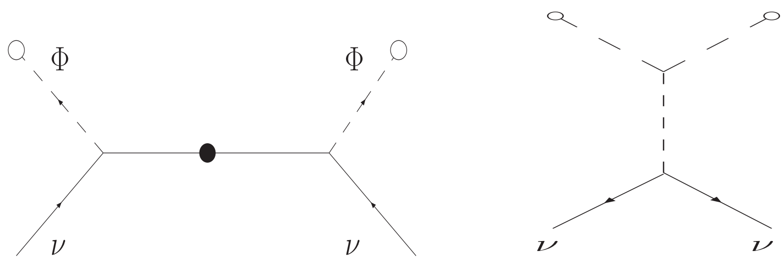

FIGURE 5. Two types of seesaw mechanism

Seesaw basics [7]

The most general seesaw is described in terms of the SM gauge structure. Most of the low energy phenomenology, such as that of neutrino oscillations, is blind to the details of the underlying seesaw theory at high energies, e. g. its gauge group, multiplet content or the nature of B-L. The full seesaw mass matrix including the SU(2) triplet (type-II) terms was first given in [7] and reads

$$\mathscr{M}_\nu = \begin{pmatrix} M_1 & D \\ D^T & M_2 \end{pmatrix}. \tag{2}$$

in the basis of "left" and "right" neutrinos ν_L and ν_L^c. Here we use the original notation of reference [7], where the "Dirac" entry is proportional to $\langle\Phi\rangle$, the M_1 comes from a triplet vev, and M_2 is a gauge singlet. The particular case $M_1 = 0$ was first mentioned in Ref. [14] and, subsequently, in [7] and [15, 18].

Note that the matrix $\mathscr{M}_\nu$ is complex, so are its Yukawa coupling sub-matrices D, M_1 and M_2, the last two symmetric, by the Pauli principle. It is diagonalized by a unitary transformation U_ν,

$$\nu_i = \sum_{a=1}^{6} (U_\nu)_{ia} n_a, \tag{3}$$

so that

$$U_\nu^T \mathscr{M}_\nu U_\nu = \mathrm{diag}(m_i, M_i). \tag{4}$$

This yields 6 mass eigenstates, including the 3 light neutrinos with masses m_i, and 3 two-component heavy leptons of masses M_i. The light neutrino mass states ν_i are given in terms of the flavour eigenstates via eq. (3). The effective light neutrino mass reads,

$$m_\nu \simeq M_1 - D{M_2}^{-1}D^T. \tag{5}$$

The smallness of light neutrino masses comes from the hierarchy $M_2 \gg D \gg M_1$. A dynamical understanding of this hierarchy is obtained in schemes where lepton number symmetry is broken spontaneously, either with gauged or ungauged B-L.

Simplest seesaw dynamics [17]

The simplest seesaw is based on the $SU(3) \otimes SU(2) \otimes U(1)$ gauge group with ungauged lepton number. The mass terms in eq. (2) are given by triplet, doublet and singlet vevs, respectively, as [17]

$$\mathscr{M}_\nu = \begin{pmatrix} Y_3 v_3 & Y_\nu \langle\Phi\rangle \\ {Y_\nu}^T \langle\Phi\rangle & Y_1 v_1 \end{pmatrix} \tag{6}$$

As already mentioned, Y_ν, Y_3 and Y_1 are complex. Neutrino masses arise either by heavy $SU(3) \otimes SU(2) \otimes U(1)$ singlet “right-handed” neutrino exchange (type I) or by the small effective triplet vev (type II), as illustrated in Fig. 5. The effective light neutrino mass is easily obtained from Eq. (5) and its diagonalization matrices containing the CP phases relevant in leptogenesis (see below) are given explicitly as a matrix perturbation series expansion in DM_2^{-1} [17].

Since lepton number is ungauged, there is a physical Goldstone boson associated with its spontaneous breakdown, the majoron. Its profile can be determined just by analysing the symmetry properties of the scalar potential (not its detailed form) which dynamically determines the vevs appearing in Eq. (6) [17]. These obey a simple hierarchy

$$v_1 \gg v_2 \gg v_3$$

with a vev seesaw relation of the type $v_3 v_1 \sim {v_2}^2$ where $v_2 \equiv \langle\Phi\rangle$ denotes the SM Higgs doublet vev, fixed by the W-boson mass. This hierarchy implies that the triplet vev $v_3 \to 0$ as the singlet vev v_1 grows and hence the type-II term is also suppressed. This model provides the first realization of seesaw that gives a dynamical understanding of the smallness of both type-I and type-II terms.

Left-right symmetric seesaw [14, 15]

This is a more symmetric (less general) version of the seesaw, where lepton number (B-L) is gauged. For example, it can be realized either in terms of $SO(10)$ or its $SU(3) \otimes SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$ subgroup [14, 15]. In $SO(10)$ each matter generation is assigned to a **16** (spinorial) so that the **16** . **16** . **10** and **16** . **126** . **16** couplings generate all entries of the seesaw matrix in Eq. (5) where Y_L and Y_R denote the Yukawas of the **126** of $SO(10)$, whose vevs $\langle\Delta_{L,R}\rangle$ give rise to the Majorana terms. They correspond to Y_1 and Y_3 of the simplest seesaw model. On the other hand Y_ν denotes the **16** . **16** . **10** Dirac Yukawa coupling. The diagonalization can be worked out as in the simplest case.

With obvious changes, e. g. $v_1 \to \langle\Delta_R\rangle$ and $v_3 \to \langle\Delta_L\rangle$, the explicit formulas for the 6×6 unitary diagonalizing matrix U given in Ref. [17] also hold.

The only important difference with respect to the previous case is the absence of the majoron, now absorbed as the longitudinal mode of the gauge boson coupled to B-L, which picks up a huge mass. The smallness of neutrino masses gets correlated to the observed maximality of parity violation in low-energy weak interactions, a connection which is as elegant as phenomenologically irrelevant, given the large value of the B-L scale required both to fit the neutrino masses, as well as to unify the gauge couplings.

Extended seesaw

In any gauge theory one can add any number of (anomaly-free) gauge singlet leptons [7]. For example, in $SO(10)$ and $E(6)$ one may add leptons outside the **16** or the **27**, respectively. Some of these extended seesaw schemes [60] are motivated by string theories [61]. New features emerge when the seesaw is realized non-minimally. Recent examples are type-III [62, 63, 64] and the $SO(10)$ seesaw mechanism with low B-L scale in Ref. [65]. For a brief review see Ref. [59].

Low-scale models

There are many models of neutrino mass where the dim-5 operator is induced from physics at low scales, TeV or less. The smallness of its strength comes then from loop and Yukawa couplings suppression and by small lepton number violating parameters that appear in its numerator, instead of its denominator. Here is an example.

Inverse seesaw [60]

It has the same mass matrix as the double seesaw model [59], except that the basic L-violating scale μ is taken very small, e. g. $\mu \ll Y_\nu\langle\phi\rangle \ll M$ [60]. As a result neutrino masses vanish as $\mu \to 0$,

$$m_\nu = \langle\Phi\rangle^2 Y_\nu^T M^{T^{-1}} \mu M^{-1} Y_\nu,$$

opposite to what happens in minimal seesaw. The entry μ may be proportional to the vev of an SU(2) singlet scalar, in which case spontaneous B-L violation leads to the existence of a majoron [66], implying a new phase transition after the electroweak transition. Since all particles are at the TeV scale, there are possibly testable phenomenological implications, including lepton flavour violation in muon and tau decays [42].

The model is “natural” in t’Hooft’s sense [67]: *“an otherwise arbitrary parameter may be taken as small when the Lagrangean symmetry increases by having it vanish”*.

Radiative models

Neutrino masses may be induced by calculable loop corrections [44, 45] as illustrated in Fig. 6. For example in the the two-loop model one has, up to a logarithmic factor,

$$\mathscr{M}_\nu \sim \lambda_0 \left(\frac{1}{16\pi^2}\right)^2 f Y_l h Y_l f^T \frac{\langle\Phi\rangle^2}{(m_k)^2} \langle\sigma\rangle \tag{7}$$

in the limit where the doubly-charged scalar k is much heavier than the singly charged one. Here l denotes a charged lepton, f and h are their Yukawa coupling matrices and Y_l denotes the SM Higgs Yukawa couplings to charged leptons. Here $\langle\sigma\rangle$ denotes an $SU(3) \otimes SU(2) \otimes U(1)$ singlet vev used in Ref. [68]. Clearly, even if the proportionality factor λ_0 is large, the neutrino mass is suppressed by the presence of a product of five small Yukawas and the appearance of the two-loop factor.

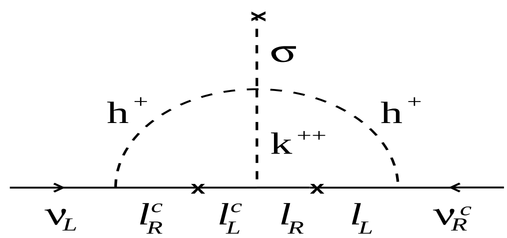

FIGURE 6. Two-loop origin for neutrino mass [45, 68].

Supersymmetry and neutrino mass

The intrinsically supersymmetric way to break lepton number is to break the so-called R parity. This may happen spontaneously, driven by a nonzero vev of an $SU(3) \otimes SU(2) \otimes U(1)$ singlet sneutrino [69, 70, 71], leading to an effective model with bilinear violation of R parity [72]. This provides the minimal way to add neutrino masses to the MSSM, we call it RMSSM [73], to stress that it serves as a reference model. Neutrino mass generation is hybrid, with one scale generated at tree level and the other induced by "calculable" radiative corrections [74]. The neutrino mass spectrum is typically "normal hierarchy"-type, with the atmospheric scale generated at the tree level and the solar mass scale arising from calculable loops, as in Fig. 7. The general form of the expression is

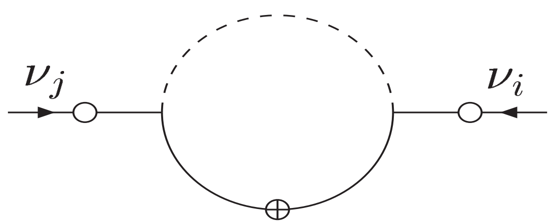

FIGURE 7. Loop-induced solar scale in RMSSM [74]; open blobs denote $\Delta L = 1$ insertions.

quite involved but the approximation

$$\mathscr{M}_\nu \sim \left(\frac{1}{16\pi^2}\right) \langle\Phi\rangle^2 \frac{A}{m_0} Y_d Y_d \tag{8}$$

(where A denotes the trilinear soft supersymmetry breaking coupling) holds in some regions of parameters.

Predicting neutrino masses and mixing

Currently five of the basic parameters of the lepton sector are probed in neutrino oscillation studies. Data points towards a well defined pattern of neutrino mixing angles, quite distinct from that of quarks, and difficult to account for in unified schemes where quarks and leptons are related. There seems to be an intriguing complementarity between quark and lepton mixing angles [75, 76, 77, 78].

There have been many papers trying to understand the values of the leptonic mixing angles from underlying symmetries. Of course, this is part of the the problem of predicting quark and lepton mixings, a defying challenge for model-builders.

Harrison, Perkins & Scott have suggested [79] that at high scales the neutrino mixing angles are given by,

$$\begin{aligned} \tan^2\theta_{\text{ATM}} &= \tan^2\theta^0_{23} = 1 \\ \sin^2\theta_{\text{Chooz}} &= \sin^2\theta^0_{13} = 0 \\ \tan^2\theta_{\text{SOL}} &= \tan^2\theta^0_{12} = 0.5. \end{aligned} \tag{9}$$

Such pattern [80] could result from some flavour symmetry. Its predictions should be corrected by renormalization group evolution [81, 82, 83].

Here I consider a specific idea to predict neutrino masses and mixing angles: that neutrino masses arise from a common seed at some “neutrino mass unification” scale M_X [84], very similar to the merging of the SM gauge coupling constants at high energies due to supersymmetry [85]. Although in its simplest form this idea is now inconsistent (at least if CP is conserved) with the observed value of the solar mixing angle θ_{12}, there is an alternative realization in terms of an A_4 flavour symmetry which is both viable and predictive [86]. Starting from three-fold degeneracy of the neutrino masses at the seesaw scale, the model predicts maximal atmospheric angle and vanishing θ_{13},

$$\theta_{23} = \pi/4 \quad \text{and} \quad \theta_{13} = 0\,.$$

Although the solar angle θ_{12} is unpredicted, one expects [1]

$$\theta_{12} = \mathscr{O}(1).$$

If CP is violated θ_{13} becomes arbitrary and the Dirac phase is maximal [88]. One can show that lepton and slepton mixings are closely related and that there must exist at least one slepton below 200 GeV, which can be produced at the LHC. The absolute Majorana neutrino mass scale $m_0 \geq 0.3$ eV ensures that the model will be probed by future cosmological tests and $\beta\beta_{0\nu}$ searches. Rates for lepton flavour violating processes $l_j \to ł_i + \gamma$ typically lie in the range of sensitivity of coming experiments, with $\text{BR}(\mu \to e\gamma) \gtrsim 10^{-15}$ and $\text{BR}(\tau \to \mu\gamma) > 10^{-9}$.

[1] There have been realizations of the A_4 symmetry that also predict the solar angle, e. g. Ref. [87].

Absolute scale of neutrino mass and $0\nu\beta\beta$

Neutrino oscillations are blind to whether neutrinos are Dirac or Majorana. Lepton number violating processes, such as $0\nu\beta\beta$ and neutrino transition electromagnetic moments [89, 90] [91, 92] probe the basic nature of neutrinos. Neutrinoless double beta decay offers the best hope. Its significance stems from the fact that, in a gauge theory, irrespective of the mechanism that induces $0\nu\beta\beta$, it necessarily implies a Majorana neutrino mass [10], as illustrated in Fig. 8. Thus it is a basic issue. Quantitative implica-

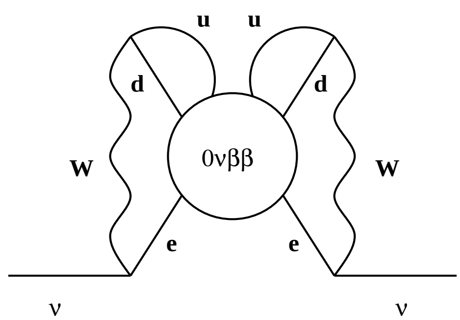

FIGURE 8. Neutrinoless double beta decay and Majorana mass are equivalent [10].

tions of the “black-box” argument are model-dependent, but the theorem itself holds in any “natural” gauge theory (for a recent discussion see [93]).

$0\nu\beta\beta$ will test absolute neutrino masses, inaccessible in neutrino oscillations, and also complement direct information that will become available from high sensitivity beta decay studies [94], as well as cosmic microwave background and large scale structure observations [95, 96, 97].

The oscillation signal implies that $0\nu\beta\beta$ must be induced by the exchange of light Majorana neutrinos, through the so-called "mass-mechanism". The corresponding amplitude is sensitive both to the absolute scale of neutrino mass, and to Majorana phases [7], neither of which can be probed in oscillations [11, 12]. Taking into account current neutrino oscillation parameters [23] and state-of-the-art nuclear matrix elements [98] one can determine the average mass parameter $\langle m_\nu \rangle$ characterizing the neutrino exchange contribution to $0\nu\beta\beta$, as shown in Fig. 10 of Ref. [59]. Models with quasi-degenerate neutrinos [86] [99] [100] give the largest $0\nu\beta\beta$ signal. For models with normal hierarchy there is in general no lower bound on $\langle m_\nu \rangle$ since there can be a destructive interference amongst the neutrino amplitudes (for an exception, see Ref. [87]; in that specific model a lower bound on $\langle m_\nu \rangle$ exists, which depends, as expected, on the value of the Majorana CP violating phase ϕ_1). In contrast, the inverted neutrino mass hierarchy implies a “lower” bound for the $0\nu\beta\beta$ amplitude.

The best current limit on $\langle m_\nu \rangle$ comes from the Heidelberg-Moscow experiment. There is also a claim made in Ref. [101] (see also [102]) which will be important to confirm or refute in future experiments. GERDA will provide an independent check of this claim [103]. SuperNEMO, CUORE, EXO, MAJORANA and possibly other experiments will further extend the sensitivity of current $0\nu\beta\beta$ searches [104].

Other phenomena

Besides oscillations and $0\nu\beta\beta$ neutrino masses may have other phenomenological manifestations. Here I summarize a few.

- **lepton flavour violation** Now that lepton flavour violation has been shown to occur in neutrino propagation it is natural to expect that it may show up elsewhere. Indeed, it is expected to occur in seesaw-type schemes of neutrino mass, either from neutral heavy lepton exchange [105, 106, 107] of through supersymmetric contributions [108, 109, 110, 111]. Note that since flavor and CP violation can occur in the massless neutrino limit, the allowed rates are unsuppressed by the smallness of neutrino masses [105, 106, 112, 113]. In the extended seesaw scheme [60] one can understand the interplay of both types of contributions. It is shown [42] that $Br(\mu \to e\gamma)$ and the nuclear $\mu^- - e^-$ conversion rates lie within planned sensitivities of future experiments such as PRISM [114].
- **TeV neutral heavy leptons** Extended seesaw models like the inverse seesaw may contain quasi-Dirac neutral heavy leptons around TeV or so, that may be directly produced at accelerators [115].
- **majoron-emitting neutrino decays** If neutrino masses arise from a spontaneous breaking of global lepton number the associated Goldstone boson (majoron) may lead to neutrino decays [17]. Although these are rather slow, they may be astrophysically relevant and lead to interesting signals [116] at underground detectors.
- **TeV new gauge boson coupled to lepton number** If neutrino masses arise from spontaneous breaking of gauged lepton number [123, 65], there will exist a light new neutral gauge boson, Z' that could be detected in searches for Drell-Yan processes at the LHC.
- **invisible Higgs boson decays** In low-scale models of neutrino mass with spontaneous breaking of global lepton number the majoron can lead to an invisible Higgs boson decays [117, 118, 119, 120].

  $$H \to JJ \tag{10}$$

  where J is the majoron. The latter is experimentally detectable as missing energy or transverse momentum associated to the Higgs [121, 122], a signal that must be taken into account when designing Higgs boson search strategies at future collider experiments. This shows that, although neutrino masses are small, the neutrino mass generation may have very important implications for the mechanism of electroweak symmetry breaking.
- **Reconstructing neutrino mixing at accelerators** Low-scale models of neutrino mass offer the tantalizing possibility of reconstructing neutrino mixing at high energy accelerators, like the LHC and the ILC. A remarkable example is provided by models where supersymmetry is the origin of neutrino mass [73]. A general feature of these models is that, unprotected by any symmetry, the lightest super-

symmetric particle (LSP) is unstable. In order to reproduce the masses indicated by current neutrino oscillation data, the LSP is expected to decay inside the detector [74] [124]. More strikingly, LSP decay properties correlate with the neutrino mixing angles. For example, if the LSP is the lightest neutralino, it should have the same decay rate into muons and taus, since the observed atmospheric angle is close to $\pi/4$ [125, 126, 127]. Such correlations hold irrespective of which supersymmetric particle is the LSP [128] and constitute a smoking gun signature of this proposal that will be tested at upcoming accelerators.

Thermal leptogenesis

It has long been noted [129] that seesaw models open an attractive possibility of accounting for the observed cosmological matter-antimatter asymmetry in the Universe through leptogenesis [130]. In this picture the decays of the heavy "right-handed" neutrinos present in the seesaw play a crucial role. These take place through diagrams in Fig. 9. In order to induce successful leptogenesis the decay must happen before the electroweak phase transition [131] and must also happen out-of-equilibrium, i. e. the decay rate must be less than the Hubble expansion rate at that epoch. Another crucial ingredient is CP violation in the lepton sector. The lepton (or B-L) asymmetry thus

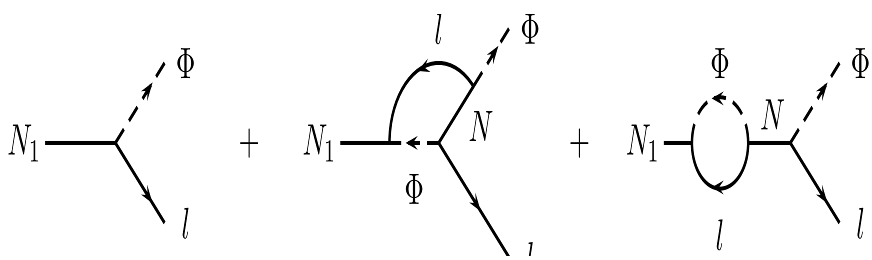

FIGURE 9. Diagrams contributing to leptogenesis.

produced then gets converted, through sphaleron processes, into the observed baryon asymmetry.

In seesaw-type schemes the high temperature needed for leptogenesis leads to an overproduction of gravitinos, which destroys the standard predictions of Big Bang Nucleosynthesis (BBN). In minimal supergravity models, with $m_{3/2} \sim 100$ GeV to 10 TeV gravitinos are not stable, decaying during or after BBN. Their rate of production can be so large that subsequent gravitino decays completely change the standard BBN scenario. To prevent such "gravitino crisis" one requires an upper bound on the reheating temperature T_R after inflation, since the abundance of gravitinos is proportional to T_R. This leads to a stringent upper bound [132], which is in conflict with the temperature required for leptogenesis, $T_R > 2 \times 10^9$ GeV [133]. One way to cure this conflict [134] is to add a small R-parity violating $\lambda_i \hat{\nu}^c{}_i \hat{H}_u \hat{H}_d$ term in the superpotential, where $\hat{\nu}^c{}_i$ are right-handed neutrino supermultiplets. One can show that in the presence of this term, the produced lepton-antilepton asymmetry can be enhanced. An alternative suggestion [135] was made in the context of extended supersymmetric seesaw schemes. It was shown in this case that leptogenesis can occur at the TeV scale through the decay of a new singlet, thereby avoiding the gravitino crisis. Washout of the asymmetry is effectively suppressed by the absence of direct couplings of the singlet to leptons.

Acknowledgments:

I thank the organizers for hospitality. This work was supported by Spanish grants FPA2005-01269, European commission RTN Contract MRTN-CT-2004-503369 and ILIAS/N6 Contract RII3-CT-2004-506222.

REFERENCES

1. Fukuda, S., et al., *Phys. Lett.*, **B539**, 179–187 (2002).
2. Ahmad, Q. R., et al., *Phys. Rev. Lett.*, **89**, 011301 (2002).
3. Kajita, T., *New J. Phys.*, **6**, 194 (2004).
4. Araki, T., et al., *Phys. Rev. Lett.*, **94**, 081801 (2004).
5. Ahn, M. H., et al., hep-ex/0606032.
6. Tagg, N. [MINOS Collaboration], eConf **C060409**, 019 (2006).
7. Schechter, J., and Valle, J. W. F., *Phys. Rev.*, **D22**, 2227 (1980).
8. Elliott, S. R., and Vogel, P., *Ann. Rev. Nucl. Part. Sci.*, **52**, 115–151 (2002).
9. Doi, M., Kotani, T., and Takasugi, E., *Prog. Theor. Phys. Suppl.*, **83**, 1 (1985).
10. Schechter, J., and Valle, J. W. F., *Phys. Rev.*, **D25**, 2951 (1982).
11. Bilenky, S. M., Hosek, J., and Petcov, S. T., *Phys. Lett.*, **B94**, 49 (1980).
12. Schechter, J., and Valle, J. W. F., *Phys. Rev.*, **D23**, 1666 (1981).
13. Doi, M., Kotani, T., Nishiura, H., Okuda, K., and Takasugi, E., *Phys. Lett.*, **B102**, 323 (1981).
14. Minkowski, P., *Phys. Lett.*, **B67**, 421 (1977).
15. Orloff, J., Lavignac, S., and Cribier, M. (2004), (Editors), written and Round table contributions in SEESAW25: Int. Conf. on the Seesaw Mechanism and the Neutrino Mass, Paris, France.
16. Weinberg, S., *Phys. Rev.*, **D22**, 1694 (1980).
17. Schechter, J., and Valle, J. W. F., *Phys. Rev.*, **D25**, 774 (1982).
18. Lazarides, G., Shafi, Q., and Wetterich, C., *Nucl. Phys.*, **B181**, 287 (1981).
19. Bahcall, J. N., and Pinsonneault, M. H., *Phys. Rev. Lett.*, **93**, 121301 (2004).
20. Honda, M., Kajita, T., Kasahara, K., and Midorikawa, S. (2004).
21. Mikheev, S. P., and Smirnov, A. Y., *Sov. J. Nucl. Phys.*, **42**, 913–917 (1985).
22. Wolfenstein, L., *Phys. Rev.*, **D17**, 2369 (1978).
23. Maltoni, M., Schwetz, T., Tortola, M. A., and Valle, J. W. F., *New J. Phys.*, **6**, 122 (2004), app. C in hep-ph/0405172 (v5) provides updated results as of June 2006; previous references given therein.
24. Bahcall, J. N., Serenelli, A. M., and Basu, S. astro-ph/0511337.
25. Aharmim, B., et al., *Phys. Rev.*, **C72**, 055502 (2005).
26. Pakvasa, S., and Valle, J. W. F. (2003), proc. of the Indian National Academy of Sciences on Neutrinos, Vol. 70A, No.1, p.189 - 222 (2004), Eds. D. Indumathi, M.V.N. Murthy and G. Rajasekaran.
27. Gonzalez-Garcia, M. et al., *Phys. Rev.*, **D63**, 033005 (2001).
28. Schechter, J., and Valle, J. W. F., *Phys. Rev.*, **D21**, 309 (1980).
29. Alsharoa, M. M., et al., *Phys. Rev. ST Accel. Beams*, **6**, 081001 (2003).
30. Apollonio, M., et al. hep-ph/0210192, CERN Yellow Report on the Neutrino Factory.
31. Albright, C., et al. hep-ex/0008064, report to the Fermilab Directorate.
32. Huber, P., Lindner, M., Rolinec, M., Schwetz, T., and Winter, W., *Phys. Rev.*, **D70**, 073014 (2004).
33. Akhmedov, E. K., Tortola, M. A., and Valle, J. W. F., *JHEP*, **05**, 057 (2004).
34. Burgess, C. P., et al., *JCAP*, **0401**, 007 (2004).
35. Burgess, C. P., Dzhalilov, N. S., Rashba, T. I., Semikoz, V. B., and Valle, J. W. F., *Mon. Not. Roy. Astron. Soc.*, **348**, 609 (2004).
36. Burgess, C., et al., *Astrophys. J.*, **588**, L65 (2003).
37. Miranda, O. G., et al., *Nucl. Phys.*, **B595**, 360–380 (2001).

38. Guzzo, M., et al., *Nucl. Phys.*, **B629**, 479–490 (2002).
39. Barranco, J., et al., *Phys. Rev.*, **D66**, 093009 (2002).
40. Miranda, O. G., Rashba, T. I., Rez, A. I., and Valle, J. W. F., *Phys. Rev.*, **D70**, 113002 (2004).
41. Valle, J. W. F., *Phys. Lett.*, **B199**, 432 (1987).
42. Deppisch, F., and Valle, J. W. F., *Phys. Rev.*, **D72**, 036001 (2005);
Deppisch, F., Kosmas, T. S., and Valle, J. W. F. *Nucl. Phys.* **B752**, 80 (2006).
43. Hall, L. J., Kostelecky, V. A., and Raby, S., *Nucl. Phys.*, **B267**, 415 (1986).
44. Zee, A., *Phys. Lett.*, **B93**, 389 (1980).
45. Babu, K. S., *Phys. Lett.*, **B203**, 132 (1988).
46. Fornengo, N., et al., *Phys. Rev.*, **D65**, 013010 (2002).
47. Huber, P., and Valle, J. W. F., *Phys. Lett.*, **B523**, 151–160 (2001).
48. Miranda, O. G., Tortola, M. A., and Valle, J. W. F., *JHEP*, **10**, 008 (2006).
49. Huber, P., Schwetz, T., and Valle, J. W. F., *Phys. Rev. Lett.*, **88**, 101804 (2002);
Phys. Rev. **D66**, 013006 (2002)
50. Barranco, J. et. al., *JHEP.*, **512**, 021 (2005).
51. Dienes, K. R., Dudas, E., and Gherghetta, T., *Nucl. Phys.*, **B557**, 25 (1999).
52. Arkani-Hamed, N., et al. *Phys. Rev.*, **D65**, 024032 (2002).
53. Faraggi, A. E., and Pospelov, M., *Phys. Lett.*, **B458**, 237–244 (1999).
54. Dvali, G. R., and Smirnov, A. Y., *Nucl. Phys.*, **B563**, 63–81 (1999).
55. Mohapatra, R. N., Nandi, S., and Perez-Lorenzana, A., *Phys. Lett.*, **B466**, 115–121 (1999).
56. Barbieri, R., Creminelli, P., and Strumia, A., *Nucl. Phys.*, **B585**, 28–44 (2000).
57. Ioannisian, A., and Valle, J. W. F., *Phys. Rev.*, **D63**, 073002 (2001).
58. de Gouvea, A., and Valle, J. W. F., *Phys. Lett.*, **B501**, 115–127 (2001).
59. Valle, J. W. F. hep-ph/0608101, review lectures at the Corfu Summer Institute on Elementary Particle Physics, September 2005 and references therein.
60. Mohapatra, R. N., and Valle, J. W. F., *Phys. Rev.*, **D34**, 1642 (1986).
61. Witten, E., *Nucl. Phys.*, **B258**, 75 (1985).
62. Akhmedov, E., Lindner, M., Schnapka, E., and Valle, J. W. F., *Phys. Rev.*, **D53**, 2752–2780 (1996);
Phys. Lett. **B368**, 270 (1996)
63. Barr, S. M., and Dorsner, I., *Phys. Lett.*, **B632**, 527–531 (2006).
64. Fukuyama, T., Ilakovac, A., Kikuchi, T., and Matsuda, K., *JHEP*, **06**, 016 (2005).
65. Malinsky, M., Romao, J. C., and Valle, J. W. F., *Phys. Rev. Lett.*, **95**, 161801 (2005).
66. Gonzalez-Garcia, M. C., and Valle, J. W. F., *Phys. Lett.*, **B216**, 360 (1989).
67. ’t Hooft, G. (1979), lecture given at Cargese Summer Inst., Cargese, France, Aug 26 - Sep 8, 1979.
68. Peltoniemi, J. T., and Valle, J. W. F., *Phys. Lett.*, **B304**, 147–151 (1993).
69. Masiero, A., and Valle, J. W. F., *Phys. Lett.*, **B251**, 273–278 (1990).
70. Romao, J. C., Santos, C. A., and Valle, J. W. F., *Phys. Lett.*, **B288**, 311–320 (1992).
71. Romao, J. C., Ioannisian, A., and Valle, J. W. F., *Phys. Rev.*, **D55**, 427–430 (1997).
72. Diaz, M. A., Romao, J. C., and Valle, J. W. F., *Nucl. Phys.*, **B524**, 23–40 (1998).
73. Hirsch, M., and Valle, J. W. F., *New J. Phys.*, **6**, 76 (2004).
74. Hirsch, M., et al., *Phys. Rev.*, **D62**, 113008 (2000), err-ibid. **D65**:119901,2002.
75. Raidal, M., *Phys. Rev. Lett.*, **93**, 161801 (2004).
76. Minakata, H., and Smirnov, A. Y., *Phy. Rev.*, **D70**, 073009 (2004).
77. Ferrandis, J., and Pakvasa, S., *Phys. Rev.*, **D71**, 033004 (2005).
78. Dighe, A., Goswami, S., and Roy, P., *Phys. Rev.*, **D73**, 071301 (2006).
79. Harrison, P. F., and Scott, W. G., *Phys. Lett.*, **B535**, 163–169 (2002).
80. Harrison, P. F., Perkins, D. H., and Scott, W. G., *Phys. Lett.*, **B530**, 167 (2002).
81. Altarelli, G., and Feruglio, F., *Nucl. Phys.*, **B720**, 64–88 (2005).
82. Hirsch, M., Ma, E., Romao, J. C., Valle, J. W. F., and Villanova del Moral, A. (2006).
83. Altarelli, G., and Feruglio, F., *New J. Phys.*, **6**, 106 (2004).
84. Chankowski, P., Ioannisian, A., Pokorski, S., and Valle, J. W. F., *Phys. Rev. Lett.*, **86**, 3488 (2001).
85. Amaldi, U., de Boer, W., and Furstenau, H., *Phys. Lett.*, **B260**, 447–455 (1991).
86. Babu, K. S., Ma, E., and Valle, J. W. F., *Phys. Lett.*, **B552**, 207–213 (2003).
87. Hirsch, M., Villanova del Moral, A., Valle, J. W. F., and Ma, E., *Phys. Rev.*, **D72**, 091301 (2005).
88. Grimus, W., and Lavoura, L., *Phys. Lett.*, **B579**, 113–122 (2004).
89. Schechter, J., and Valle, J. W. F., *Phys. Rev.*, **D24**, 1883 (1981), err. D25, 283 (1982).

90. Wolfenstein, L., *Phys. Lett.*, **B107**, 77 (1981).
91. Pal, P. B., and Wolfenstein, L., *Phys. Rev.*, **D25**, 766 (1982).
92. Kayser, B., *Phys. Rev.*, **D26**, 1662 (1982).
93. Hirsch, M., Kovalenko, S., and Schmidt, I., *Phys. Lett.*, **B642**, 106 (2006).
94. Drexlin, G., *Nucl. Phys. Proc. Suppl.*, **145**, 263–267 (2005).
95. Lesgourgues, J., and Pastor, S., *Phys. Rep.*, **429**, 307–379 (2006).
96. Hannestad, S. hep-ph/0602058.
97. Fogli, G. L., et al., *Phys. Rev.*, **D70**, 113003 (2004).
98. Bilenky, S. M., Faessler, A., and Simkovic, F., *Phys. Rev.*, **D70**, 033003 (2004).
99. Caldwell, D. O., and Mohapatra, R. N., *Phys. Rev.*, **D48**, 3259–3263 (1993).
100. Ioannisian, A., and Valle, J. W. F., *Phys. Lett.*, **B332**, 93–99 (1994).
101. Klapdor-Kleingrothaus, H. V., et al., *Phys. Lett.*, **B586**, 198–212 (2004).
102. Aalseth, C. E., et al., *Mod. Phys. Lett.*, **A17**, 1475–1478 (2002).
103. Aalseth, C. E., et al., *Phys. Rev.*, **D65**, 092007 (2002).
104. Saakyan, R., Nones, C., Tomei, C., and Zuber, K. (2006).
105. Bernabeu, J., et al., *Phys. Lett.*, **B187**, 303 (1987).
106. Gonzalez-Garcia, M. C., and Valle, J. W. F., *Mod. Phys. Lett.*, **A7**, 477–488 (1992).
107. Ilakovac, A., and Pilaftsis, A., *Nucl. Phys.*, **B437**, 491 (1995).
108. Hall, L. J., Kostelecky, V. A., and Raby, S., *Nucl. Phys.*, **B267**, 415 (1986).
109. Borzumati, F., and Masiero, A., *Phys. Rev. Lett.*, **57**, 961 (1986).
110. Casas, J. A., and Ibarra, A., *Nucl. Phys.*, **B618**, 171–204 (2001).
111. Antusch, S., Arganda, E., Herrero, M. J., and Teixeira, A. hep-ph/0607263.
112. Branco, G. C., Rebelo, M. N., and Valle, J. W. F., *Phys. Lett.*, **B225**, 385 (1989).
113. Rius, N., and Valle, J. W. F., *Phys. Lett.*, **B246**, 249–255 (1990).
114. Kuno, Y., *AIP Conf. Proc.*, **542**, 220–225 (2000).
115. Dittmar, M., et al., *Nucl. Phys.*, **B332**, 1 (1990); DELPHI, P. Abreu *et al.*, *Z. Phys.* **C74**, 57 (1997).
116. Kachelriess, M., Tomas, R., and Valle, J. W. F., *Phys. Rev.*, **D62**, 023004 (2000).
117. Joshipura, A. S., and Valle, J. W. F., *Nucl. Phys.*, **B397**, 105–122 (1993); for an early paper see Shrock, R. E. and Suzuki, M. *Phys. Lett.* **B110**, 250 (1982).
118. Romao, J. C., de Campos, F., and Valle, J. W. F., *Phys. Lett.*, **B292**, 329–336 (1992).
119. Hirsch, M., et al., *Phys. Rev.*, **D70**, 073012 (2004); *Phys. Rev.*, **D73**, 055007 (2006).
120. Bazzocchi, F., and Valle, J. W. F. hep-ph/0609093.
121. de Campos, F., Eboli, O. J. P., Rosiek, J., and Valle, J. W. F., *Phys. Rev.*, **D55**, 1316–1325 (1997).
122. Abdallah, J., et al., *Eur. Phys. J.*, **C32**, 475–492 (2004).
123. Valle, J. W. F., *Phys. Lett.*, **B196**, 157 (1987).
124. de Campos, F., et al., *Phys. Rev.*, **D71**, 075001 (2005).
125. Porod, W., Hirsch, M., Romao, J., and Valle, J. W. F., *Phys. Rev.*, **D63**, 115004 (2001).
126. Romao, J. C., et al., *Phys. Rev.*, **D61**, 071703 (2000).
127. Mukhopadhyaya, B., Roy, S., and Vissani, F., *Phys. Lett.*, **B443**, 191–195 (1998).
128. Hirsch, M., and Porod, W., *Phys. Rev.*, **D68**, 115007 (2003).
129. Fukugita, M., and Yanagida, T., *Phys. Lett.*, **B174**, 45 (1986).
130. Buchmuller, W., Peccei, R. D., and Yanagida, T., *Ann. Rev. Nucl. Part. Sci.*, **55**, 311–355 (2005).
131. Kuzmin, V. A., Rubakov, V. A., and Shaposhnikov, M. E., *Phys. Lett.*, **B155**, 36 (1985).
132. Kawasaki, M., Kohri, K., and Moroi, T., *Phys. Rev.*, **D71**, 083502 (2005).
133. Buchmuller, W., Di Bari, P., and Plumacher, M. *Annals Phys.* **315**, 305 (2005)
134. Farzan, Y., and Valle, J. W. F., *Phys. Rev. Lett.*, **96**, 011601 (2006).
135. Hirsch, M., Malinsky, M., Romao, J. C., Sarkar, U., and Valle, J. W. F. hep-ph/0608006